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# Speculations on ALEPH's Dijet Enhancement

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## Abstract

We interpret the dijet enhancement reported by the ALEPH collaboration in the process  $e^+e^- \rightarrow 4$  jets as being due to the production of a pair of bottom squarks, followed by their R-parity violating decays into pairs of light quarks. Constraints on this speculative interpretation are examined. Some of the consequences of our hypothesis are drawn within the context of softly broken supersymmetry.

Recently the ALEPH collaboration presented a preliminary analysis of about  $6 \text{ pb}^{-1}$  of data collected at  $\sqrt{s} = 130 - 136 \text{ GeV}$  at LEP [1]. Although cross sections for standard processes appear to be consistent with expectations, they reported some peculiar 4-jet events in their data for which there is no canonical interpretation. What ALEPH sees is an excess of 4-jet events (14 observed, 7.1 expected), with 8 of these events clustered at a dijet “sum mass” of about 110 GeV. This dijet sum mass is arrived at by combining together the 2 dijet masses in the events which have the smallest mass difference between them. Assuming that what is seen is a signal for something, it corresponds to a cross section of  $2.5 \pm 1 \text{ pb}$ . This “signal” is quite distinct from what is expected from QCD 4-jet events, where the event distribution is essentially flat in the dijet sum mass. Indeed, in the sum mass bin from 102-116 GeV where the 8 events are clustered, one expects only 1.35 events from QCD.

ALEPH makes no claims about this signal and it could well be a statistical fluctuation which will disappear as more data is collected. Nevertheless, since the dijet mass difference in the analysis is restricted to be below 20 GeV, it is tempting to speculate that the ALEPH excess is due to the production of a pair of particles of mass of about 55 GeV which then each decay into pairs of jets. In fact, already Farrar [2] has suggested that the ALEPH events are associated with the pair production of squarks, which subsequently decay into a quark and a light gluino. Although the light gluino scenario is interesting, here we wish to speculate in another direction. We also want to associate the ALEPH signal with the pair production of squarks. However, in contrast to Farrar, we suggest that what is produced is only the  $\tilde{b}_L$  and that the dijets result from the R-parity violating decay of this squark to pairs of light quarks. In what follows we will try to justify this speculation and draw some of its consequences.

For a 55 GeV squark at LEP 1.5 energies, the cross section for  $e^+e^- \rightarrow \tilde{b}_L \tilde{b}_L^*$  is about 1 pb. Thus the size of the signal is of the right magnitude. Although the electroweak symmetry does not permit an R-parity violating trilinear term involving the  $(t, b)_L$  doublet in the superpotential, there is a possible  $\Delta B = 1$  R-parity violating term involving  $b_R$  [3]. Since the  $\tilde{b}_L$  squarks mix with the  $\tilde{b}_R$  squarks, these latter couplings (if they are present) allow for the decay of  $\tilde{b}_L$  into light quarks. One must check, however, that this decay does not run afoul of any of the existing constraints on the trilinear R-parity violating couplings.

There are nine possible  $\Delta B = 1$  trilinear R-parity violating couplings in the superpotential involving right-handed isosinglet quarks

$$W = \lambda_{ijk} D_R^i D_R^j U_R^k, \quad (1)$$

since  $\lambda_{ijk} = -\lambda_{jik}$ . The most stringent bound on these couplings is that on  $\lambda_{dsu}$ , which comes from the process  $NN \rightarrow KKX$ , where  $N$  is a nucleon and  $K$  is a particle with unit strangeness. In Ref. [4], it was shown that one could have a bound as strong as  $|\lambda_{dsu}| < 10^{-7}$ , but with large hadronic uncertainties and significant model dependence. Six of the couplings  $\lambda_{ijk}$  will involve a  $b_R$ :  $\lambda_{bdu}; \lambda_{bsu}; \lambda_{bdc}; \lambda_{bsc}; \lambda_{bdt}; \lambda_{bst}$ . The last two couplings are not relevant here since the top is too heavy to be produced in the decay of a 55 GeV  $\tilde{b}_L$ . However, assuming that  $\tilde{b}_L$  is 55 GeV, one must make sure that

$$\lambda_{bdt}; \lambda_{bst} \ll e$$

to preserve the decay  $t \rightarrow Wb$  as the main top decay mode. To our knowledge, there are no stringent bounds on  $\lambda_{bsu}$ ,  $\lambda_{bdc}$  and  $\lambda_{bsc}$ . However,  $\lambda_{bdu}$  gives a contribution to  $n - \bar{n}$  oscillations and the experimental limit on the neutron oscillation lifetime [5] of  $\tau > 1.2 \times 10^8$  sec bounds this parameter to be below about  $10^{-5}$ <sup>1</sup>. As long as some of the couplings  $\lambda_{bij}$  are of order  $10^{-5}$  or larger, the decay of the produced  $\tilde{b}_L$  into light quarks will not produce a displaced vertex<sup>2</sup>. Thus, the existence of at least one R-parity violating coupling of sufficient strength among  $\lambda_{bsu}$ ,  $\lambda_{bdc}$  and  $\lambda_{bsc}$  makes the suggested scenario phenomenologically viable. It could be argued that we have introduced unnatural hierarchies in the couplings  $\lambda_{ijk}$  by requiring that we evade certain bounds while still retaining one coupling of sufficient strength. However, it is worth pointing out that even within the minimal standard model one has sizable hierarchies in the Yukawa sector, such as  $m_u/m_t \sim 10^{-5}$ . We should note that, at this stage, one could have also imagined that the ALEPH events came from the pair production of the lightest stop, followed by their R-parity violating decays into light quarks. For this to be viable, however, one would have to imagine that the coupling  $\lambda_{sdt}$  dominated over both  $\lambda_{bst}$  and  $\lambda_{bdt}$  to account for the lack of decays with a  $b$ -jet in the final state in the ALEPH 4 jet sample [1].

To proceed, however, we must still check whether our suggestion is theoretically sound. For the decay  $\tilde{b}_L \rightarrow 2\text{-jets}$  to dominate, it is important that the  $\tilde{b}_L$  be the LSP. In particular, all neutralinos should be heavier than the  $\tilde{b}_L$ . Otherwise the decay  $\tilde{b}_L \rightarrow b\chi_1^0$  is likely to be dominant, since the relevant coupling is of  $O(e)$ . We have associated the ALEPH enhancement to the production of a  $\tilde{b}_L$  because this squark, along with the stops, has a mass that is sensitive to the large top Yukawa coupling. Thus, as is well known [6], even starting with universal SUSY-breaking scalar masses at the GUT scale, the  $\tilde{b}_L$  has the possibility of obtaining a nonuniversal mass. In fact, assuming universal soft breaking masses for the scalars and the gauginos and  $\tan\beta \sim 2$ , one finds

$$m_{\tilde{d}_L}^2 \simeq m_0^2 + 6.8m_{1/2}^2, \quad (2)$$

while

$$m_{\tilde{b}_L}^2 \simeq 0.51m_0^2 + 5.5m_{1/2}^2. \quad (3)$$

One sees from the above that it is possible to have the  $\tilde{b}_L$  mass be smaller than that of the other down squarks, provided the universal scalar mass  $m_0$  dominates over  $m_{1/2}$ . However, because we want  $m_{\tilde{b}_L} \simeq 55$  GeV then both  $m_0$  and, particularly, the universal gaugino mass  $m_{1/2}$  must be quite small. This, in general, is unacceptable since it leads to one or more very light neutralinos in the spectrum. Neutralinos in this mass range are excluded experimentally by LEP [7]. But more importantly, the presence of neutralinos lighter than the  $\tilde{b}_L$  would alter the decays of these squarks in an undesirable way. Indeed, in the presence of such light neutralinos, the weak decay  $\tilde{b}_L \rightarrow \chi_1^0 b$  followed by the R-parity violating decay  $\chi_1^0 \rightarrow 3\text{ jets}$  would give a different experimental signature from that suggested by the ALEPH data.

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<sup>1</sup>The actual bound depends in detail on the supersymmetric spectrum, as discussed by Goity and Sher [4].

<sup>2</sup>In fact, the decay of the  $\tilde{b}_L$  via the R-parity violating operator (1) can only occur in the presence of some (typically small)  $\tilde{b}_L$ - $\tilde{b}_R$  mixing. Hence we need  $\lambda_{bij} \sin\theta_b \geq 10^{-5}$ , where  $\sin\theta_b \sim 5 \times 10^{-2}$  is the sine of the  $\tilde{b}_L$ - $\tilde{b}_R$  mixing angle. However, the combination  $\lambda_{bij} \sin\theta_b$  also cannot be too large, for this would lead to  $Z \rightarrow qq'\tilde{b}_L$  decays at an unacceptable level.

The presence of light neutralinos, however, is a feature of the particular pattern of SUSY breaking which one assumes at the GUT scale. The principal contributor to the  $m_{1/2}$  piece of  $m_{\tilde{b}_L}$  in Eq. (3) is the gluino component. However, the neutralino masses are sensitively dependent on the SUSY breaking masses one gives to the **electroweak gauginos**, but weakly dependent on the gluino mass. If the electroweak gaugino masses are taken to be different from the gluino mass, it is possible to make the neutralinos sufficiently heavy by having the  $SU(2) \times U(1)$  gaugino masses  $m_1$  and  $m_2$  heavy. This, per se, should not affect terribly the  $b_L$  mass. However, to be sure one must examine anew the spectrum of supersymmetric excitations one gets in the case where the soft breaking of supersymmetry involves nonuniversal gaugino masses,  $m_1 \neq m_2 \neq m_3$ <sup>3</sup>.

We present below the results of a study of a minimal supergravity model with unequal soft supersymmetry breaking mass parameters in the gaugino sector, but where one still has a common SUSY breaking mass for all the scalars. To constrain the model further, we also assume that the electroweak symmetry is broken radiatively [6]. Even though a common mass is assumed for all the scalars at the GUT scale, these masses evolve to different values at low energy as a result of radiative effects. For the squarks of the first two generations and for the sleptons, the effective masses at low energy are sensitive functions of the gaugino masses and the evolution of the coupling constants. The masses for the third generation squarks depend, in addition, on the top Yukawa coupling and its evolution. Furthermore, for the mass squared of the  $\tilde{t}_L$  and  $\tilde{t}_R$  one cannot neglect the SUSY-preserving contribution of  $m_t^2$ , due to the large top mass.

Solving the renormalization group equations for the squark and slepton masses [6], one obtains the following approximate formulas for the light squarks and the sleptons:

$$\begin{aligned}
m_{\tilde{e}_L}^2(t) &= m_0^2 - 0.27 \cos 2\beta M_Z^2 + x_1(t)m_1^2 + x_2(t)m_2^2 \\
m_{\tilde{\nu}_L}^2(t) &= m_0^2 + 0.5 \cos 2\beta M_Z^2 + x_1(t)m_1^2 + x_2(t)m_2^2 \\
m_{\tilde{e}_R}^2(t) &= m_0^2 - 0.23 \cos 2\beta M_Z^2 + 4x_1(t)m_1^2 \\
m_{\tilde{u}_L}^2(t) &= m_0^2 + 0.35 \cos 2\beta M_Z^2 + \frac{1}{9}x_1(t)m_1^2 + x_2(t)m_2^2 + x_3(t)m_3^2 \\
m_{\tilde{d}_L}^2(t) &= m_0^2 - 0.42 \cos 2\beta M_Z^2 + \frac{1}{9}x_1(t)m_1^2 + x_2(t)m_2^2 + x_3(t)m_3^2 \\
m_{\tilde{u}_R}^2(t) &= m_0^2 + 0.15 \cos 2\beta M_Z^2 + \frac{16}{9}x_1(t)m_1^2 + x_3(t)m_3^2 \\
m_{\tilde{d}_R}^2(t) &= m_0^2 - 0.08 \cos 2\beta M_Z^2 + \frac{4}{9}x_1(t)m_1^2 + x_3(t)m_3^2 .
\end{aligned}$$

Here  $t = \ln \frac{M_X^2}{M_Z^2}$  and the  $x_i(t)$  are functions that depend on the running of the standard model coupling constants. Taking  $M_X = 2 \times 10^{16}$  GeV and  $\alpha_{GUT} = \frac{1}{24.3}$  [9] one has, approximately,

$$x_1(t) = 0.038; \quad x_2(t) = 0.49; \quad x_3(t) = 6.30.$$

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<sup>3</sup> Nonuniversal gaugino masses have been advocated recently by Roszkowski and Shifman[8] in a different context.

For the third generation squarks, except  $\tilde{b}_R$ , there are additional contributions due to the top Yukawa coupling. One finds:

$$\begin{aligned} m_{\tilde{b}_R}^2(t) &= m_{\tilde{d}_R}^2(t) \\ m_{\tilde{b}_L}^2(t) &= m_{\tilde{d}_L}^2(t) + h_0(t)m_0^2 + h_{ij}(t)m_i m_j \\ m_{\tilde{t}_R}^2(t) &= m_{\tilde{u}_R}^2(t) + 2h_0(t)m_0^2 + 2h_{ij}(t)m_i m_j + m_t^2 \\ m_{\tilde{t}_L}^2(t) &= m_{\tilde{u}_L}^2(t) + h_0(t)m_0^2 + h_{ij}(t)m_i m_j + m_t^2 \end{aligned} .$$

Here the functions  $h_0(t)$  and  $h_{ij}(t)$  depend on the evolution of the top Yukawa coupling. For example, using as input <sup>4</sup>  $\tan \beta = 2.4$  and  $m_t \simeq 180$  GeV, one has

$$\begin{aligned} h_0(t) &= -0.494 \\ h(t) &= \begin{pmatrix} -0.0154 & -0.0004 & -0.0024 \\ -0.0004 & -0.1214 & -0.0203 \\ -0.0024 & -0.0203 & -1.0793 \end{pmatrix} . \end{aligned}$$

Because  $h_0(t)$  is negative, the scalar mass contribution for  $\tilde{b}_L$ ,  $\tilde{t}_L$  and  $\tilde{t}_R$  is smaller than that for the other squarks. Furthermore, because  $m_{\tilde{d}_L}^2$  is very weakly dependent on  $m_1^2$  and  $h_{11} < 0$ , one sees that large values of the  $U(1)$ -gaugino mass will further reduce the  $\tilde{b}_L$  mass provided that the contributions of the other two gaugino masses are relatively contained. This amounts to a considerably fine-tuned cancelation of the  $U(1)$  gaugino contribution against the contributions of the other gauginos. A large value for  $m_1$  has the desired property of raising the neutralino masses, securing the role of  $\tilde{b}_L$  as the LSP. Taking, for example,  $m_0 = 20$  GeV and

$$m_1 = 1900 \text{ GeV}; \quad m_2 = 160 \text{ GeV}; \quad m_3 = 80 \text{ GeV}$$

gives  $m_{\tilde{b}_L} \simeq 56$  GeV, which is more than 50 GeV below the next two lightest squarks ( $m_{\tilde{t}_1} \simeq 115$  GeV and  $m_{\tilde{u}_L} \simeq 260$  GeV). <sup>5</sup>

Once all the SUSY-breaking parameters  $m_i$  are fixed (along with the SUSY-preserving mass parameter  $\mu$ , which is then a function of  $\tan \beta$ ), it is straightforward to deduce the gluino mass and the neutralino and chargino mass spectra. For the parameters detailed above, one finds:

$$m_{\tilde{g}} = 230 \text{ GeV},$$

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<sup>4</sup>It is not necessary to assume a value for  $\tan \beta$  if one assumes [10] that the top Yukawa coupling at the top mass is that which corresponds to the IR fixed point. Although the value for  $\tan \beta$  we shall use gives a  $Y_t(m_t)$  close to the fixed point value, we shall take  $\tan \beta$  as a free parameter, and fix the top Yukawa coupling using the known top mass  $180 \pm 12$  GeV [11].

<sup>5</sup>Here  $\tilde{t}_1$  is the lighter of the two stops and the detailed value of its mass depends on the mixing between  $\tilde{t}_L$  and  $\tilde{t}_R$ . This in turns depends on the value of the trilinear SUSY-breaking parameter  $A_0$  at the GUT scale and on the supersymmetric mass parameter  $\mu$ . We have taken  $A_0 = 0$  and have fixed  $\mu$  from the minimization condition which must be satisfied to obtain  $SU(2) \times U(1)$  breaking [10]. For this example one has  $\mu = 302$  GeV.

while

$$m_{\chi_1^0} = 106 \text{ GeV}; \quad m_{\chi_2^0} = 305 \text{ GeV}; \quad m_{\chi_3^0} = 327 \text{ GeV}; \quad m_{\chi_4^0} = 788 \text{ GeV}$$

and

$$m_{\chi_1^\pm} = 106 \text{ GeV}; \quad m_{\chi_2^\pm} = 332 \text{ GeV}.$$

For these same parameters the *tree level* Higgs masses are

$$m_{H_1} = 64 \text{ GeV}; \quad m_{H_2} = 532 \text{ GeV}; \quad m_A = 528 \text{ GeV}; \quad m_{H^+} = 534 \text{ GeV}.$$

To get an idea of how “fine tuned” the above mass spectrum is, we present in Fig.1 a scatter plot of minimal neutralino and chargino masses obtained by varying the input parameters  $m_i$  and  $\tan \beta$ , but requiring still that  $m_{\tilde{b}_L} \simeq 55 \text{ GeV}$ .

By lifting the assumption of universal gaugino masses at the GUT scale, we have succeeded in producing a spectrum of SUSY particles which exhibits the essential features necessary to account for the ALEPH “signal”: the LSP is a bottom squark of mass  $\sim 55 \text{ GeV}$  which decays to pairs of jets via  $R$ -parity violating couplings. However, to establish the phenomenological viability of this scenario, we must still examine various indirect constraints on the model. Two such constraints are provided by the  $\rho$ -parameter [12] and SUSY-mediated flavor changing neutral currents. A further set of constraints can be obtained from analyzing new top quark decay modes that are present in this scenario.

The supersymmetric contributions to the  $\rho$  parameter have been studied previously in Refs. [13, 14]. It has been shown that the extra Higgs particles and the gauginos of the SUSY standard model give a negligible contribution to  $\rho$ . However, non-degenerate  $SU(2)_L$  squark doublets can give a large contribution to  $\rho$ , in complete analogy with the similar result for quarks in the minimal standard model. In particular, for the spectrum described above, the large splitting between the top and bottom squarks can give a sizable contribution to  $\rho$ : writing  $\rho = 1 + \delta\rho_{\text{MSM}} + \delta\rho_{\text{SUSY}}$ , we have [13, 14]

$$\begin{aligned} \delta\rho_{\text{SUSY}} = & \frac{3\alpha}{8\pi M_W^2 \sin^2 \theta_W} \bigg[ \\ & \cos^2 \theta_t \left( \cos^2 \theta_b f(m_{\tilde{t}_1}, m_{\tilde{b}_1}) + \sin^2 \theta_b f(m_{\tilde{t}_1}, m_{\tilde{b}_2}) \right) \\ & + \sin^2 \theta_t \left( \cos^2 \theta_b f(m_{\tilde{t}_2}, m_{\tilde{b}_1}) + \sin^2 \theta_b f(m_{\tilde{t}_2}, m_{\tilde{b}_2}) \right) \\ & - \cos^2 \theta_t \sin^2 \theta_b f(m_{\tilde{t}_1}, m_{\tilde{t}_2}) - \cos^2 \theta_b \sin^2 \theta_b f(m_{\tilde{b}_1}, m_{\tilde{b}_2}) \bigg], \end{aligned}$$

where  $\theta_t$  is the mixing angle between  $\tilde{t}_L$  and  $\tilde{t}_1$ ,  $\theta_b$  is the mixing angle between  $\tilde{b}_L$  and  $\tilde{b}_1$ , and  $f(m_1, m_2)$  is given by

$$f(m_1, m_2) = \frac{m_1^2 + m_2^2}{2} - \frac{m_1^2 m_2^2}{m_1^2 - m_2^2} \log \frac{m_1^2}{m_2^2}.$$

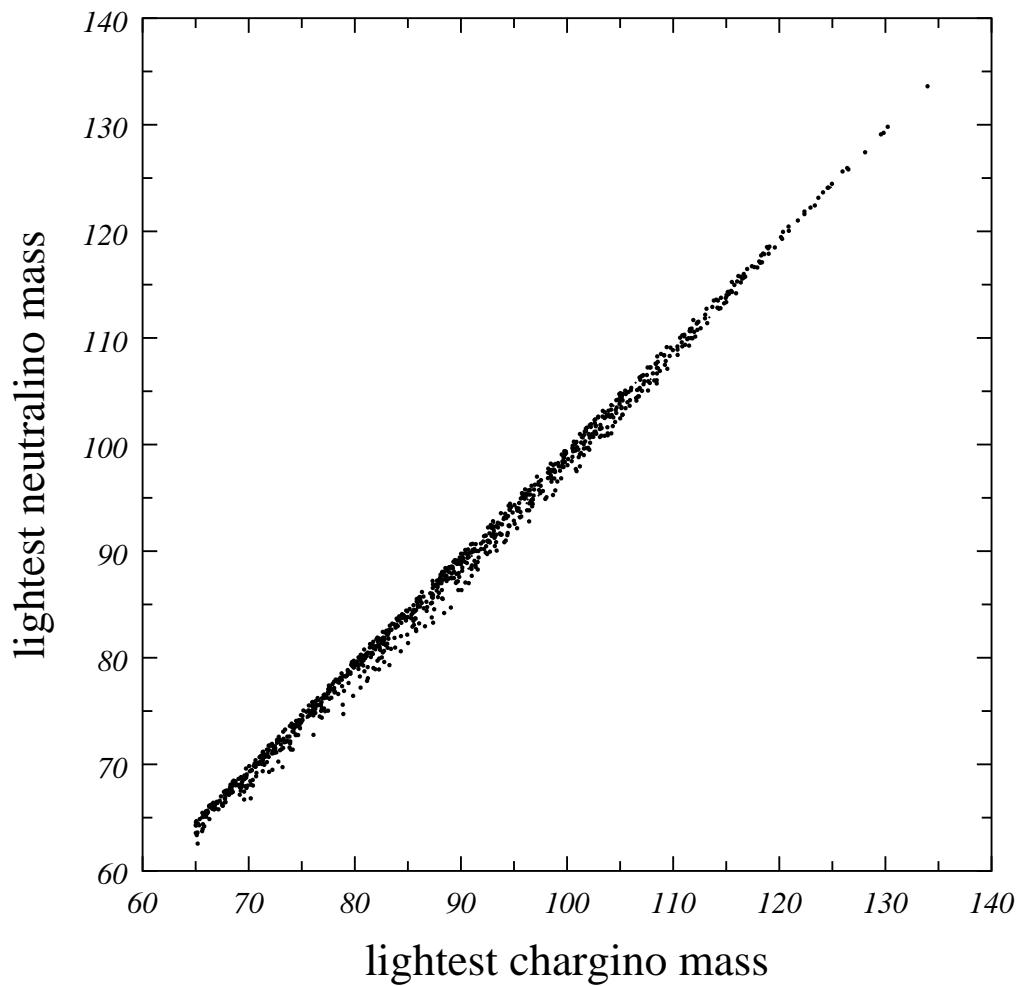


Figure 1: Lightest chargino and neutralino masses obtained for a sample of 1000 spectra with  $m_{\tilde{b}_L} = 55 \pm 5$  GeV.

For a top mass of  $180 \pm 12$  GeV, a fit to the data yields a liberal<sup>6</sup> bound on  $\delta\rho_{\text{SUSY}}$  of 0.004. Evaluating  $\delta\rho_{\text{SUSY}}$  for the spectrum given above, we find that the squarks give a significant, but not unacceptable, contribution to the  $\rho$  parameter:  $\delta\rho_{\text{SUSY}} = 0.0028$ .

Supersymmetric particles with the mass spectrum considered here can also have effects on flavor changing neutral currents in the  $B$  system. In particular, it is known that supersymmetry can give dangerously large enhancements of the  $b \rightarrow s\gamma$  decay rate. The branching ratio for charmless radiative  $B$  decays has been measured by CLEO [16]; at the 95% confidence level, they have reported

$$1 \times 10^{-4} < \text{BR}(B \rightarrow X_s\gamma) < 4 \times 10^{-4}.$$

The contribution of supersymmetric particles to this decay rate can be computed using the results of Refs. [17, 18, 19, 20]. There it was shown that

$$\frac{\text{BR}(B \rightarrow X_s\gamma)}{\text{BR}(B \rightarrow X_c e \bar{\nu})} = \frac{|V_{ts}^* V_{tb}|^2}{|V_{cb}^2|} \frac{6\alpha_{\text{QED}}}{\pi g(m_c/m_b)} |C_7(\mu)|^2$$

where  $g(m_c/m_b)$  is the phase space factor for the semileptonic decay, and  $C_7(\mu)$  is the coefficient of the flavor-changing operator

$$O_7 = \frac{e}{4\pi^2} m_b \bar{s}_L \sigma_{\mu\nu} b_R F^{\mu\nu}$$

evaluated at a scale  $\mu \sim m_b$ . The QCD evolution of  $C_7(\mu)$  has been computed in Refs. [17, 18], and the contributions of supersymmetric particles to the relevant Wilson coefficients can be found in Ref. [21]. In the leading logarithmic approximation, we have [19]

$$C_7(\mu) \simeq \eta^{\frac{16}{23}} C_7(M_W) + \frac{8}{3} \left( \eta^{\frac{14}{23}} - \eta^{\frac{16}{23}} \right) C_8(M_W) + \sum_{i=1}^8 h_i \eta^{a_i},$$

where  $C_8$  is the coefficient of the chromo-magnetic moment operator

$$O_8 = \frac{g_s}{4\pi^2} m_b \bar{s}_L \sigma_{\mu\nu} T^a b_R G_a^{\mu\nu},$$

and  $\eta = \alpha_s(M_W)/\alpha_s(\mu)$ . The coefficients  $h_i$  and  $a_i$  are pure numbers independent of any model parameters, and can be found in Ref. [19].

The coefficients  $C_7(M_W)$  and  $C_8(M_W)$  can be found in Ref. [21]. Normalizing them appropriately, the standard model contributions are given by

$$C_7^{W^\pm}(M_W) = \frac{3x^3 - 2x^2}{4(x-1)^4} \log x + \frac{-8x^3 - 5x^2 + 7x}{24(x-1)^3},$$

and

$$C_8^{W^\pm}(M_W) = \frac{-3x^2}{4(x-1)^4} \log x + \frac{-x^3 + 5x^2 + 2x}{8(x-1)^3},$$

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<sup>6</sup>By this we mean that this is the maximum allowed value of  $\delta\rho$  when its correlation with the Peskin-Takeuchi  $S$  parameter [15] is taken into account, and  $S$  is allowed to vary such that the 90% confidence level limit on  $\delta\rho_{\text{SUSY}}$  is maximized. If we fix  $S = 0$ , the bound decreases to 0.003.

where  $x = m_t^2/M_W^2$ . In the present case, we anticipate that the charged Higgs, the charginos, and the gluino will give the bulk of the supersymmetric contribution. The charged Higgs has long been known to give a sizable enhancement of radiative  $B$  decay rates [17]. The charginos can further enhance or suppress the rate [21]. The presence of relatively small gluino and  $\tilde{b}_L$  masses in the spectrum given above indicate that the gluino may also give a sizable contribution. The contributions of the charged Higgs and the charginos to  $C_{7,8}(M_W)$  may be found, for instance, in Ref. [22]. For the gluino, we have [21]

$$C_7^{\tilde{g}}(M_W) = \frac{\alpha_s \sin^2 \theta_W}{\alpha} \frac{U_{b\tilde{b}_L} U_{s\tilde{b}_L}^*}{V_{tb} V_{ts}^*} \frac{M_W^2}{m_{\tilde{b}_L}^2} \left[ \frac{4z^2}{9(z-1)^4} \log z + \frac{-4z^2 - 10z + 2}{27(z-1)^3} \right],$$

and

$$C_8^{\tilde{g}}(M_W) = \frac{\alpha_s \sin^2 \theta_W}{2\alpha} \frac{U_{b\tilde{b}_L} U_{s\tilde{b}_L}^*}{V_{tb} V_{ts}^*} \frac{M_W^2}{m_{\tilde{b}_L}^2} \left[ \frac{-9z + z^2}{3(z-1)^4} \log z + \frac{-11z^2 + 40z + 19}{18(z-1)^3} \right],$$

where  $z = m_{\tilde{g}}^2/m_{\tilde{b}_L}^2$ , and  $U_{q\tilde{q}_L}$  is the flavor mixing matrix that appears in the coupling of a quark to a gluino and a left-handed squark. We have neglected a small additional contribution that involves  $\tilde{b}_L$ – $\tilde{b}_R$  mixing. Similar “super-CKM” matrices appear in the couplings of the charginos to bottom quarks and top squarks. Given the large non-degeneracy of the squarks in the above spectrum, it would be reasonable to expect that these “super CKM” matrices may deviate from the standard model CKM matrix. Hence in the following we will retain these CKM elements to make explicit the dependence on these angles.

Evaluating  $C_{7,8}$  for the spectrum given above ( $m_{\tilde{b}_L} = 56$  GeV and  $m_{\tilde{g}} = 230$  GeV) and including the contribution of the charged Higgs and the charginos, we find an estimate for  $\Gamma(b \rightarrow s\gamma)$ :

$$\frac{\Gamma(b \rightarrow s\gamma)[\text{SUSY}]}{\Gamma(b \rightarrow s\gamma)[\text{MSM}]} \simeq |1 + 0.18 + 0.09\sigma - 0.49\lambda|^2$$

where the four terms come from  $W$ , charged Higgs, gluino, and chargino loops. The contributions of *all* of the squarks have been included and expressed in terms of third generation mixing angles using CKM unitarity. The quantities  $\sigma$  and  $\lambda$  are given by

$$\sigma = \frac{U_{b\tilde{b}_L} U_{s\tilde{b}_L}^*}{V_{tb} V_{ts}^*}$$

and

$$\lambda = \frac{\tilde{V}_{tb} \tilde{V}_{ts}^*}{V_{tb} V_{ts}^*}.$$

The matrix  $\tilde{V}$  is the super CKM matrix describing couplings of charginos to up-type squarks and down-type quarks. We see that for  $\lambda \sim 1$  (a reasonable value), the charginos interfere destructively with the  $W$  and charged Higgs, reducing the rate. The possibility of such a phenomenon has been noted in Ref. [22]. Given the standard model estimate [20]

$$\text{BR}(B \rightarrow X_s \gamma)[\text{MSM}] \simeq 1.9 \pm 0.5 \times 10^{-4},$$

we conclude that the rate is compatible with the CLEO determination.

The presence of light supersymmetric particles in the spectrum allows the top quark to decay in other modes besides the standard  $t \rightarrow Wb$  mode. For the model spectrum discussed above, the dominant non standard decay is  $t \rightarrow \tilde{b}_L \chi_1^+$ <sup>7</sup>. The branching ratio for this mode is of  $O(30\%)$  which, although sizable, is probably acceptable given the experimental uncertainty in the top cross section [11]. Because the dominant chargino decay is  $\chi_1^+ \rightarrow \tilde{b}_L^* c$ , with the  $\tilde{b}_L$  decaying into dijets, the final state for these non standard top decays will contain five jets. This is not a particularly easy signal to dig out. Nevertheless, if some top tagging can eventually be implemented at the Tevatron, looking for multijet decays of the accompanying  $\bar{t}$  may be the best way to dig out the  $\tilde{b}_L$  in hadronic interactions<sup>8</sup>.

It is quite likely that, in the end, the ALEPH dijet enhancement will prove to be just a statistical fluctuation, making the scenario discussed here moot. Even if this were to turn out to be the case, we believe that some of our disquisitions may continue to prove useful. Three lessons which emerge from our speculations stand out in particular. First, significant deviations from the spectrum predicted by minimal versions [9, 10] of softly broken supersymmetry can occur as a result of some simple changes in the underlying assumptions (e.g. having non-universal gaugino masses). Second, a low mass sbottom (or a low mass stop) should not be unexpected, given the large top Yukawa coupling. Third, although there is a natural prejudice against R-parity violating couplings, their presence at some level is by no means ruled out. If present, such couplings totally alter the “standard” signals of supersymmetry.

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<sup>7</sup> Another allowed decay is  $t \rightarrow \tilde{t}_1 \tilde{b}_L b$ , but this decay is strongly kinematically suppressed. There may also be R-parity violating top decay modes into  $bd$  and  $bs$ , but the relevant branching fractions for these modes depend on the unknown couplings  $\lambda_{bit}$ .

<sup>8</sup> Although  $\tilde{b}_L - \tilde{b}_L$  pairs are copiously produced at the Tevatron, because of the large QCD background it appears very difficult to detect them with an ALEPH-like analysis.

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